B Physics and Hadronic Matrix Elements [*]

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In a mini workshop on B Physics which were held at Dalian, China, a number of lectures was given. Around 40 participants joined the workshop. In this workshop, people discussed the heavy quark effective field theory, light cone QCD sum rules and light cone wave functions. Much discussion was devoted to the currently popular QCD factorization and perturbative QCD approaches in the hadronic B decays. The physics of CP violation and physics beyond standard model were also discussed.

I. OUTLINE

The standard model (SM) has been well tested in the gauge sector, the unsolved and unclear problems in the SM are mainly concerned in the Yukawa sector which is strongly related to flavor physics, such as: origin and mechanism of CP violation, origin of the quark and lepton masses as well as their mixing. It involves thirteen parameters whose origin are all unknown. Therefore, precisely extracting those parameters and testing CP violation mechanism as well as probing new physics become hot topics in flavor physics. In fact, flavor physics has already indicated the existence of new physics. After forty years of discovery for indirect CP violation, direct CP violation in kaon decays has well been established. Theoretical predictions [1] and experimental measurements [2, 3] are now consistent each other. Exclusive semi-leptonic and inclusive B decays play a crucial role for extracting two important parameters V_{cb} and V_{ub} in the CKM matrix elements. Rare B decays and direct CP violations are also of great importance in determining weak phase angles of the unitarity triangle and testing the Kobayashi-Maskawa (KM) mechanism [4] in SM as well as probing new physics.

The currently running B-Factories at KEK and SLAC provide more and more experimental data on B physics. The future LHCb, BTeV and super-B factory will give much more data. More Precise Experimental Data and more precision theoretical prediction will encourage precision test of SM at B sector.

In this workshop (next section), Y.L. Wu talks about how the heavy quark effective field theory provides a powerful tool for studying beauty physics, what is the implication of beauty physics. Some attentions are also paid to the more precise extraction of V_{cb} and V_{ub} and large direct CP violation in charmless B decays.

In section III, T. Huang gives a brief review on the light-cone wave function(LCWF) in QCD. The definition of the light-cone wave function in QCD and distribution amplitude and its asymptotic form are reviewed. General properties of the light-cone wave function and higher helicity components are discussed. C.F. Qiao discusses about the derivation of B meson wave function. X.H. Wu talks about the twist 3 wave function of pion in QCD sum rules. The $1/m_b$ power suppressed effects are discussed by Z.H. Li.

Y.D. Yang introduces the QCD factorization approach and discusses the application in hadronic B decays in section IV. K.T. Chao intensively studies the B decays to charmonium final states in QCD factorization in section V. X.Q. Yu and C.D. Lu present the formalism and application of perturbative QCD approach, which is based on k_T factorization in section VI. Intensive discussions are induced for the comparison and underlining theory of these two approaches.

The CKM angle determination is one of the main talks for the B factories. Z.J. Xiao discusses a number of channels which are useful for these angles α, β and γ in section VII. In section VIII, C.H. Chang discuss the exact solution of Bethe-Salpeter (BS) equation.

The new physics picture is always a hot topic. X.G. He and C.S. Huang give the SUSY contributions in $B \to K^* \gamma$ decay and $B \to \phi K_S$ and $B \to \eta' K_S$ decays in section IX. These decays are more sensitive to new physics than some others. Z.J. Xiao makes a systematic study for the new physics contributions to the charmless two-body hadronic

decays of B and B_s meson induced by the neutral- and charged-Higgs penguin diagrams in model III: the third type of the two-Higgs-doublet models and technicolor models.

II. LCQCD/HQEFT FROM FULL QCD

by Y.L. Wu

Heavy quark effective field theory of QCD which was first explored in [5] and detailed developed recently in a serious papers [6, 7, 8, 9, 10, 11, 12, 13, 14, 15, 16, 17] provides a premising and systematic tool in correctly evaluating the hadronic matrix elements of heavy quarks and extracting the CKM matrix elements V_{cb} and V_{ub} from B decays via heavy quark expansion (HQE). HQEFT is a theoretical framework derived directly from QCD but explicitly displays the heavy quark symmetry in the infinite mass limit $m_Q \to \infty$ and symmetry breaking corrections for finite mass case in the real world. In fact, HQEFT has been shown to be as a large component QCD [5, 15]. At the leading order, it coincides with the usual heavy quark effective theory(HQET) [18] which is constructed based on the heavy quark symmetry (HQS) [19] in the infinite mass limit [20]. The differences between HQEFT of QCD and HQET arise from the sub-leading terms in the $1/m_Q$ expansion. This is because in the construction of HQET the particle and antiparticle components were separately treated based on the assumption that the particle number and antiparticle number are conserved in the effective Lagrangian [21]. However, such an assumption is only valid in the infinite mass limit. Note that the particle number and antiparticle number are always conserved in the transition matrix, which is independent of heavy quark limit. Therefore the particle-antiparticle coupling terms which correspond to the pair creation and annihilation terms in full QCD were inappropriately dropped away in the HQET though they are suppressed by $1/m_Q$ but only vanish in infinite mass limit. In HQEFT of QCD [5, 15], all contributions of the field components, large and small, 'particle' and 'antiparticle', have carefully been treated in the effective Lagrangian, so that the resulting effective Lagrangian form the basis for a "complete" effective field theory of heavy quarks. It is seen that the particle-antiparticle coupling terms are suppressed by $1/m_Q$ and they are decoupled in the infinite mass limits. Their physics effects at sub-leading order have been found to be significant in some cases [6, 7, 8, 9, 10, 11, 12, 13, 14]. Therefore, to consider the finite quark mass correction precisely, it is necessary to include the contributions from the components of the anti-quark fields. Note that the Wilson coefficient functions describe the perturbative effects in full QCD, HQEFT of QCD treats the non-perturbation effects below the energy scale m_b. So that the anti-quark effects in HQEFT of QCD cannot be attributed to the Wilson coefficient functions. As the anti-quark effects appear from the sub-leading order of $1/m_Q$ expansion, it should not surprise that at the leading order the resulting anomalous dimensions from both full QCD and HQEFT of QCD are all the same.

A complete description for the theoretical framework of LCQCD/HQEFT is going to be presented in a longer paper with the following contents:

1. Introduction
2. Effective Lagrangian of LCQCD/HQEFT from QCD
3. Lorentz Invariance of LCQCD/HQEFT
4. Quantization of LCQCD/HQEFT
4.1 Quantum Generators of Poincare Group
4.2 Anti-commutations and Super selection Rule
4.3 Hilbert and Fock Space of LCQCD/HQEFT
4.4 Propagator in LCQCD/HQEFT
4.5 Discrete Symmetries in LCQCD/HQEFT
5. Basic Framework of LCQCD/HQEFT
5.1 Feynman Rules in LCQCD/HQEFT
5.2 Finite Mass Corrections
5.3 Trivialization of Gluon Couplings and Decouple Theorem
5.4 Renormalization in LCQCD/HQEFT and Wilson Loop
5.5 Current Operators in LCQCD/HQEFT
6. New Symmetries in LCQCD/HQEFT
7. Demonstration for the Trace Formula of Transition Matrix Elements and Universal Isgur-Wise Function
8. Simple applications
Hereafter is a summary of Important Effects in LCQCD/HQEFT:

- Automatically reparameterization invariance $v \to v'$ with $v^2 = v'^2 = 1$.

- Automatically with vanishing $1/m_Q$ corrections at Zero-Recoil (both for $B \to D^* l \nu_l$ & $B \to D l \nu_l$)

- Largely reduce the numbers of form factors at higher Order of $1/m_Q$ and more relations are obtained

- Nicely keep scaling law of decay constants

$f_H = \frac{F}{\sqrt{M_H}}[1 + O(\frac{1}{m_Q})]$ ($\sim 10\%$ for b , $\sim 30\%$ for c)

The explicit forms and results were found to be as follows

$$f_M = \frac{F}{\sqrt{m_M}}\{1 + \frac{1}{m_Q}(g_1 + 2d_M g_2)\}$$

$$\frac{f_V m_V^{1/2}}{f_P m_P^{1/2}} = 1 - \frac{8}{m_Q}g_2$$

$$F = 0.38 \pm 0.06 \text{ GeV}^{3/2}; \quad g_1 = 0.46 \pm 0.12 \text{ GeV}; \quad g_2 = -0.06 \pm 0.02 \text{ GeV}$$

The numerical results are

$$f_B(m_b) = 0.196 \pm 0.044 \text{ GeV}, \qquad f_{B^*}(m_b) = 0.206 \pm 0.039 \text{ GeV},$$
$$f_D(m_c) = 0.298 \pm 0.109 \text{ GeV}, \qquad f_{D^*}(m_c) = 0.354 \pm 0.090 \text{ GeV}.$$

- Reliably understand the lifetime differences of bottom hadrons B_d, B_s, Λ_b.

$$\frac{\tau(\Lambda_b)}{\tau(B^0)} = 0.79 \pm 0.01$$

$$\frac{\tau(B^-)}{\tau(B_d)} = 1.08 \pm 0.05 \,, \qquad \frac{\tau(B_s)}{\tau(B_d)} = 0.96 \pm 0.06$$

$$B_{SL}(B^0) = (10.51 \pm 0.37), \qquad B_{SL}(\Lambda_b) = (11.13 \pm 0.26)\%$$

$$n_c(B^0) = 1.20 \pm 0.03, \qquad B_\tau(B^0) = (2.6 \pm 0.1)\%$$

- Consistently extraction of V_{cb} and V_{ub}

$$|V_{cb}| = 0.0395 \pm 0.0011_{\text{exp}} \pm 0.0019_{\text{th}} \quad \text{from} \quad B \to D^* l\nu, \quad O(1/m_Q^2)$$
$$|V_{cb}| = 0.0434 \pm 0.0041_{\text{exp}} \pm 0.0020_{\text{th}} \quad \text{from} \quad B \to Dl\nu, \quad O(1/m_Q^2)$$
$$|V_{cb}| = 0.0394 \pm 0.0010_{\text{exp}} \pm 0.0014_{\text{th}} \quad \text{from} \quad B \to X_c l\nu, \quad O(1/m_Q^2)$$
$$|V_{cb}| = 0.0402 \pm 0.0014_{\text{exp}} \pm 0.0017_{\text{th}} \quad (Average)$$
$$|V_{ub}|_{LO} = (3.4 \pm 0.5_{\text{exp}} \pm 0.5_{\text{th}}) \times 10^{-3} \quad \text{from} \quad B \to \pi l\nu,$$
$$|V_{ub}|_{LO} = (3.7 \pm 0.6_{\text{exp}} \pm 0.7_{\text{th}}) \times 10^{-3} \quad \text{from} \quad B \to \rho l\nu,$$
$$|V_{ub}|_{NLO} = (3.2 \pm 0.5_{\text{exp}} \pm 0.2_{\text{th}}) \times 10^{-3} \quad \text{from} \quad B \to \pi l\nu,$$
$$|V_{ub}| = (3.48 \pm 0.62_{\text{exp}} \pm 0.11_{\text{th}}) \times 10^{-3} \quad \text{from} \quad B \to X_u l\nu \quad O(1/m_Q^2)$$

III. LIGHT-CONE SUM RULES AND LIGHT-CONE WAVE FUNCTION IN QCD

by T. Huang, Z.H. Li, C.F. Qiao and X.H. Wu

A brief review on the light-cone wave function in QCD is presented.It consists of the following nine sections: 1) the definition of the light-cone wave function in QCD; 2) distribution amplitude(DA) and its asymptotic form; 3) general properties of the light-cone wave function; 4) Melosh transformation and higher helicity components; 5) DA moments in the QCD sum rules; 6) model approach to LCWF; 7) Conformal expansion; 8) Twist-3 distribution amplitude of the pion; 9) perspective on the light-cone wave function.

As well known,the light cone formalism provides a convenient framework for the relativistic description of hadrons in terms of quark and gluon degrees of freedom, and the application of perturbative QCD (PQCD) to exclusive processes has mainly been developed in this formalism [24].In the PQCD theory, the hadronic distribution amplitudes and structures which enter exclusive and inclusive processes via the factorization theorems at high momentum transfer can be determined by the hadronic wave functions, and they are the underlying links between hadronic phenomena in QCD at large distances (nonperturbative) and small distance (perturbative). A light-cone (LC) wave function is

defined at the LC time $\tau \equiv x^+ = x^0 + x^3$ in physical light cone gauge $A^+ = A^0 + A^3 = 0$, which is conjugate to the LC Hamiltonian $H_{LC} \equiv P^- = P^0 - P^3$ [25].A stationary solution $|\Psi(\tau)\rangle$ has a LC Fock state expansion and the LC wave functions, $\Psi_n(x_i, \mathbf{k}_{\perp i}, \lambda_i)$, are the amplitudes to find n partons (quarks, anti-quarks and gluons) with momenta k_i in a pion of momentum P. Hence it is important to develop methods which specify the light-cone wave functions of the hadron.

The distribution amplitude $\Phi(x, Q)$ for the leading twist satisfies a QCD evolution equation [24].As Q^2 goes to infinity,$\Phi(x, Q)$ is dominated mainly by the first term in the expansion,i.e for the twist-2 DA of the meson,we have $\Phi_M(x, Q) \to a_M^{(0)} x(1 - x)$.However the solution of the evolution equation can be determined completely only as if the initial condition of $\Phi(x, Q_0)$ is given.

From the relation between the wave functions and measurable quantities we can get the general properties of the hadronic wave functions [25]. For example,in the pion case two important constraints on the valence state wave function $\Psi_{q\bar{q}}(x, \mathbf{k}_\perp)$ (for the $\lambda_1 + \lambda_2 = 0$ components) have been derived from $\pi \to \mu\nu$ and $\pi^0 \to \gamma\gamma$ decay amplitude. The quark transverse momentum of the valence state in the pion $\langle \mathbf{k}_\perp^2 \rangle_{q\bar{q}}$ should be larger than $\langle \mathbf{k}_\perp^2 \rangle_\pi$. $P_{q\bar{q}}$ is the probability of finding the $q\bar{q}$ Fock state in the pion, and it should be less than unity. the charged pion radius for the lowest valence quark state $\langle r_{\pi^+}^2 \rangle^{q\bar{q}}$ should be lower than the experimental value for $\langle r_{\pi^+}^2 \rangle$, where the latter one should contain all the contributions of the higher Fock stats.

In order to obtain the LC spin wave function of the pion, we transform the ordinary instant-form SU(6) quark model wave function into a LC one. The instant-form spin states $|J, s\rangle_T$ and the LC-form spin states $|J, \lambda\rangle_F$ are related by $|J, \lambda\rangle_F = \sum_s U_{s\lambda}^J |J, s\rangle_T$. This rotation is called the Melosh rotation for spin-1/2 particles. Applying the Melosh rotation we can obtain the spin wave function of the pion in the infinite-momentum frame [27]. There are two higher helicity components ($\lambda_1 + \lambda_2 = \pm 1$) in the expression of the LC spin wave function besides the ordinary helicity components ($\lambda_1 + \lambda_2 = 0$). These higher helicity components will make contributions to the exclusive processes although they have the power suppression behavior.

Applying the QCD sum rules to the distribution amplitude,we can get the first three moments of the initial DA,$\Phi(x, Q_0)$ [26]. From these moments we can not determine the initial $\Phi(x, Q_0)$ and only know that the initial $\Phi(x, Q_0)$ is different from the asymptotic form.

For the space wave function,one can take the Brodsky-Huang-Lepage (BHL) prescription which is obtained by connecting the equal-time wave function in the rest frame and the wave function in the infinite momentum frame [25]. We obtain the LC space wave function of the pion, $\varphi_{\text{BHL}} = A \exp[-\frac{\mathbf{k}_\perp^2 + m^2}{8\beta^2 x(1-x)}]$ [25]. This form provides a more reasonable k_T dependence in the space wave function and is consistent with the experimental data [28].

Conformal expansion allows one to separate transverse and longitudinal variables in the distribution amplitude.one can expand DA in terms of Gegenbauer polynomials $C_n^{1/2}$ and $C_n^{3/2}$ according to its conformal spin [29]. This symmetry helps us to make a model for the light-cone wave function of the hadron.

Recently, we do not use the equation of motion for the quarks inside the pion and apply the QCD sum rules to calculate the pion twist-3 DA $\Phi_p(\xi)$ and get its approximate expression and its expansion coefficients can be determined by its three moments approximately [30]. Based on the moment calculation we suggest a model for the twist-3 wave function with k_T dependence [31].

In Summary, at present although we can not solve light-cone wave function of the hadron completely due to its non-perturbative dynamics, one has got much progress by studying its general properties,symmetry, its moments from non-perturbative calculation and phenomenological model. These studies will contribute to a better understanding on the light-cone wave functions and approach their realistic solutions.

A. Twist-3 Distribute Amplitude of the Pion in QCD Sum Rules

We apply the background field method to calculate the moments of the pion two-particle twist-3 distribution amplitude (DA) $\phi_p(\xi)$ in QCD sum rules. The QCD sum rules were used to study the leading-twist distribution amplitude of the pion at the first time in [26]. The pion twist-3 distribution amplitudes have been studied in Ref. [32, 33] in the chiral limit, based on the techniques of nonlocal product expansion and conformal expansion and including corrections in the meson-mass. However they employ the equations of motion of on-shell quarks in the meson to get two relations among two-particles twist-3 distribution amplitudes of the pion, $\phi_p(\xi)$ and $\phi_\sigma(\xi)$, and three-particle twist-3 distribution amplitude $\phi_{3\pi}(\alpha_i)$ of the pion. The question is whether one can use the equation of motion of the quark inside the meson since the quarks are not on-shell. In our paper [30], we do not apply the quark equation of motion and calculate the moments of the twist-3 distribution amplitudes $\phi_p(\xi)$ directly in QCD sum rules approach.

For calculating the moments of $\phi_p(\xi)$, we take as usual the two-point correlation functions:

$$i \int d^4x e^{iq \cdot x} \langle 0 \mid T \left\{ j_5^{(2n)}(x,z) j_5^{(0)\dagger}(0) \right\} \mid 0 \rangle \equiv (z \cdot q)^{2n} I_p^{2n,0}(q^2),$$

where the currents are defined as

$$j_5^{(2n)}(x,z) = \bar{d}(x)\gamma_5 (iz \cdot \overleftrightarrow{D})^{2n} u(x)$$

with $\overleftrightarrow{D} = \overleftrightarrow{\partial_\mu} - i2g A_\mu^a T^a$. To get the sum rules for moments, we employ dispersion relation for $I_p^{2n,0}(Q^2)$ and isolate the pole term of the lowest pseudoscalar pion. With condensates up to dimension-6 and the perturbative contribution part to the lowest order, we obtain the sum rule for moments of $\phi_p(\xi)$:

$$\langle \xi_p^{2n} \rangle \cdot \langle \xi_p^0 \rangle = \frac{M^4}{(m_0^p)^2 f_\pi^2} e^{m_\pi^2/M^2} \left[\frac{3}{8\pi^2} \frac{1}{2n+1} \left(1 - (1 + \frac{s_\pi}{M^2})e^{-\frac{s_\pi}{M^2}} \right) \right.$$
$$\left. -(2n-1)\frac{\bar{m}\langle \bar{\psi}\psi \rangle}{M^4} + \frac{2n+3}{24} \frac{\langle \frac{\alpha_s}{\pi} G^2 \rangle}{M^4} - \frac{16\pi}{81} (21 + 8n(n+1)) \frac{\langle \sqrt{\alpha_s}\bar{\psi}\psi \rangle^2}{M^6} \right]$$

with $\bar{m} = (m_u + m_d)/2$, and M is a Borel parameter. For the condensate parameters, we take as usual. And the definition of the moments are given by

$$\langle \xi_p^{2n} \rangle = \frac{1}{2} \int_{-1}^{1} \xi^{2n} \phi_p(\xi)d\xi.$$

The key point to avoid the concept of on-shell equations of motion is that the introduced m_0^p appears in the sum rule of moments. And they can be obtained through requiring normalization of the zeroth moments($\langle \xi_p^0 \rangle = 1$). The advantage is that our results do not rely on the validness of the equation of motion and at the same time provide a parameter which can be compared with the corresponding one used in phenomenological analysis. With these points kept in mind, one can show that the results for the first three moments are($M^2 = 1.5 - 2 \text{ GeV}^2$):

$s_\pi(\text{GeV}^2)$	$m_0^p(\text{GeV})$	$\langle \xi_p^2 \rangle$	$\langle \xi_p^4 \rangle$
1.7	$1.24 - 1.36$	$0.340 - 0.356$	$0.167 - 0.210,$
1.6	$1.19 - 1.30$	$0.340 - 0.359$	$0.164 - 0.211,$
1.5	$1.14 - 1.24$	$0.341 - 0.361$	$0.160 - 0.212.$

It can be seen that, the parameter m_0^p is smaller than $m_\pi^2/(m_u + m_d) \simeq 1.78 - 3.92$ GeV which was used in, e.g., [32] when the equation of motion of on-shell quarks are employed. We take the first three moments into consideration in our analysis. Expanding twist-3 distribution amplitudes $\phi_p(\xi)$ of the pion in Gegenbauer polynomials to first three terms, one can obtain its approximate form,

$$\phi_p(\xi) = 1 + 0.137\, C_2^{1/2}(\xi) - 0.721\, C_4^{1/2}(\xi).$$

B. B-Meson Wave Functions

The development of the heavy quark effective theory(HQET) has led to much progress in the theoretical understanding of the properties of hadrons. With the progress in both theory and experiment, nowadays, B physics becomes one of the most active research areas in high energy physics. Many B meson exclusive decay processes turn out to be calculable systematically in the frameworks of newly developed factorization formalisms, the so-called PQCD approach [34] and QCD Factorization approach [35]. In all of these calculations based on the factorization approaches, the light-cone distribution amplitudes of the participating mesons, which express the nonperturbative long-distance contributions in the factorized amplitudes, play an important role in making reliable predictions.

By definition, the light-cone distribution amplitudes are given by the light-cone wave functions at zero transverse separation of the constituents. Following [36] it is

$$\langle 0|\bar{q}(z)\Gamma h_v(0)|\bar{B}(p)\rangle = -\frac{if_B M}{2}\text{Tr}\left[\gamma_5 \Gamma \frac{1+\not v}{2} \left\{ \tilde{\phi}_+(t) - \not z \frac{\tilde{\phi}_+(t) - \tilde{\phi}_-(t)}{2t} \right\} \right]. \tag{1}$$

Eq. (1) is the most general parametrization of the two-particle light-cone matrix element compatible with Lorentz invariance and the heavy-quark limit. The three-particle matrix element can be generally expressed as [37]:

$$\langle 0|\bar{q}(z)\, g G_{\mu\nu}(uz)\, z^\nu\, \Gamma\, h_v(0)|\bar{B}(p)\rangle = \frac{1}{2}\, f_B M\, \mathrm{Tr}\left[\gamma_5\, \Gamma\, \frac{1+\not{v}}{2}\right.$$

$$\left.\left\{(v_\mu\, \not{z} - t\,\gamma_\mu)\,\left(\tilde{\Psi}_A(t,u) - \tilde{\Psi}_V(t,u)\right) - i\,\sigma_{\mu\nu}z^\nu\,\tilde{\Psi}_V(t,u) - z_\mu\,\tilde{X}_A(t,u) + \frac{z_\mu}{t}\,\not{z}\tilde{Y}_A(t,u)\right\}\right]. \tag{2}$$

By virtue of the equations of motion, one can get a set of relations between distribution amplitudes. In the approximation that the three-particle amplitudes are set to be zero, the Wandzura-Wilczek approximation, the system of the obtained differential equations is simplified, and the solutions for $\tilde{\phi}_\pm$ can be obtained. In the momentum space, the analytic solution can be expressed explicitly as [37]

$$\psi_+(\omega, \boldsymbol{k}_T) = \frac{\omega}{2\pi\bar{\Lambda}^2}\theta(\omega)\theta(2\bar{\Lambda}-\omega)\delta\left(\boldsymbol{k}_T^2 - \omega(2\bar{\Lambda}-\omega)\right) \tag{3}$$

$$\psi_-(\omega, \boldsymbol{k}_T) = \frac{2\bar{\Lambda}-\omega}{2\pi\bar{\Lambda}^2}\theta(\omega)\theta(2\bar{\Lambda}-\omega)\delta\left(\boldsymbol{k}_T^2 - \omega(2\bar{\Lambda}-\omega)\right) \tag{4}$$

The results (3) and (4) give exact description of the valence Fock components of the B meson wave functions in the heavy-quark limit, and represent their transverse momentum dependence explicitly. These results show that the dynamics within the two-particle Fock states is determined solely in terms of a single nonperturbative parameter $\bar{\Lambda}$.

It should be noted that the Wandzura-Wilczek approximation is not equivalent to the free field approximation. The leading Fock-states, which correspond to the twist-2 contribution in the case of light meson wave functions, carries the effect from the QCD interaction.

There has been indication that, in the heavy-light quark systems, the higher Fock states could play important roles even in the leading twist level [37]. This would suggest that the shape of the wave functions as function of momenta and their quantitative role in the phenomenological applications would be modified when including the higher Fock states. However, the direct applications of our wave functions to phenomenological processes indicate that the Wandzura-Wilczek approximation does not induce obvious conflicts with experimental data [38]

In contrast to many model assumptions, which show strong damping at large $|\boldsymbol{z}_T|$ as $\sim \exp\left(-K^2\boldsymbol{z}_T^2/2\right)$, our wave functions (3) and (4) have slow-damping with oscillatory behavior as $\psi_\pm(\omega, -\boldsymbol{z}_T^2) \sim \cos(|\boldsymbol{z}_T|\sqrt{\omega(2\bar{\Lambda}-\omega)} - \pi/4)/\sqrt{|\boldsymbol{z}_T|}$. On the other hand, heavy-quark symmetry guarantees that the solution in our work provides complete description of the light-cone valence Fock wave functions for the B^* mesons and also for the D, D^* mesons in the heavy-quark limit. Moreover, the investigations on beyond Wandzura-Wilczek limit are in progress.

C. $1/m_b$ power-suppressed effects in two-body B decays

A careful investigation of $1/m_b$ power suppressed effects is crucial for a good understanding of two-body hadronic decays of B mesons. It is instructive to systematically investigate the $1/m_b$ suppressed effects from either soft or hard gluon exchanges, in addition to annihilation topology. Focusing on penguin-dominated $B \to K\pi$ and color-suppressed $\bar{B}^0 \longrightarrow D^0\pi^0$ decays, we estimate such power-suppressed corrections which appear possibly in the framework of QCD light-cone sum rules (LCSR).

In the former case, the soft exchanges could be induced by the emission topology or by the chromo-magnetic diploe operator O_{8g} and penguin topology, while the hard gluon exchanges may occur through the O_{8g} operator or penguin contractions. The behaviors of such contributions in the heavy quark limit $m_b \to \infty$ comply with the following power counting: (1) The soft effects due to the emission topology are of $\mathcal{O}(1/m_b)$. (2) The penguin diagrams make no contributions up to $1/m_b^2$ order. (3) The O_{8g} operator supplies a hard effect of $\mathcal{O}(1/m_b)$ and a soft effect of $\mathcal{O}(1/m_b^2)$. Despite being formally suppressed by $\mathcal{O}(1/m_b^2)$ compared with the leading-order factorizable amplitudes, the soft correction from O_{8g} is free of α_s suppression and numerically comparable with the $\mathcal{O}(\alpha_s)$ hard part. (4) The part with quark condensate has a chiral enhancement factor r_χ^π via the PCAC relation, suffering formally from $\mathcal{O}(\alpha_s)$ and $\mathcal{O}(1/m_b)$ double suppression but numerically turning out to be a large effect. However, it has a counterpart in the QCD factorization and thus is not included in our calculation to avoid a double counting. For the weak phase $\gamma(= \mathrm{Im}V_{ub}^*)$ ranging from $40°$ to $80°$, we find that because of the annihilation and soft and hard exchanges a numerical increase of $(20-30)\%$ is expected in the branching ratios up to $\mathcal{O}(1/m_b^2)$ corrections; the resultant soft and hard corrections are less important than the annihilation contributions and amount only to a level of 10%. Possible sources of uncertainties are discussed in some details.

For the case of $\bar{B}^0 \longrightarrow D^0\pi^0$, naive factorization assumption breaks down and non-factorizable contributions are leading. Then the $1/m_b$ power-suppressed soft exchange between the emitted heavy-light quark pair and $B\pi$ system is expected to be much more important than in the former case. The resulting correction to the decay amplitude is found to be comparable numerically with the corresponding factorizable piece, estimated at about $(50 - 110)\%$ of the latter, and renders the relevant parameter a_2 receive a positive number correction. This indicates that such power-suppressed effect would be crucial to our phenomenological understanding of the $\bar{B}^0 \longrightarrow D^0\pi^0$ decay.

IV. QCD IMPROVED FACTORIZATION APPROACH FOR B DECAYS

by Y.D. Yang

It is of great interest and importance to investigate the decays of B mesons to charmless final states to study the weak interactions and CP violation. In past years, we have witnessed many experimental and theoretical progresses in the study of B physics with the observations of many B charmless decay processes and improvements of the theory of B decays.

Most of the theoretical studies of B decays to pseudoscalar and vector final states are based on the popular *Naive Factorization* approach [39]. As it was pointed out years ago in Ref.[40], the dominant contribution in B decays comes from the so-called *Feynman mechanism*, where the energetic quark created in the weak decay picks up the soft spectator softly and carries nearly all of the final-state meson's momentum. It is also shown that Pion form factor in QCD at intermediate energy scale is dominated by Feynman mechanism [41, 42, 43]. ¿From this point, we can understand why the naive factorization approach have worked well for B and D decays, and the many existing predictions for B decays based on naive factorization and spectator ansatz do have taken in the dominant physics effects although there are shortcomings. However, with the many new data available from CLEO and an abundance of data to arrive within few years from the B factories BaBar and Belle, it is demanded highly to go beyond the naive factorization approach.

Recently, Beneke *et.al.*, have formed an interesting QCD factorization formula for B exclusive non-leptonic decays [35, 44]. The factorization formula incorporates elements of the naive factorization approach(as leading contribution) and the hard-scattering approach(as sub-leading corrections), which allows us to calculate systematically radiative(sub-leading non-factorizable) corrections to naive factorization for B exclusive non-leptonic decays. An important product of the formula is that the strong final-state interaction phases are calculable from the first principle which arise from the hard-scattering kernel and hence process dependent. The strong phases are very important for studying CP violation in B decays. Detail proofs and arguments could be found in [35].

The amplitude of B decays to two light mesons, say M_1 and M_2, is obtained through the hadronic matrix element $\langle M_1(p_1)M_2(p_2)|\mathcal{O}_i|B(p)\rangle$, here M_1 denotes the final meson that picks up the light spectator quark in the B meson, and M_2 is the another meson which is composed of the quarks produced from the weak decay point of b quark. Since the quark pair, forming M_2, is ejected from the decay point of b quark carrying the large energy of order of m_b, soft gluons with the momentum of order of Λ_{QCD} decouple from it at leading order of Λ_{QCD}/m_b in the heavy quark limit. As a consequence any interaction between the quarks of M_2 and the quarks out of M_2 is hard at leading power in the heavy quark expansion. On the other hand, the light spectator quark carries the momentum of the order of Λ_{QCD}, and is softly transferred into M_1 unless it undergoes a hard interaction. Any soft interaction between the spectator quark and other constituents in B and M_1 can be put into the transition form factor of $B \to M_1$. The non-factorizable contribution to $B \to M_1M_2$ can be calculated through the diagrams in Fig.1.

It is worth to note that the shortcomings in the "generalized factorization" are resolved in this framework. Non-factorizable effects are calculated in a rigorous way here instead of being parameterized by effective color number. Since the hard scattering kernels are convoluted with the light cone DAs of the mesons, gluon virtuality $k^2 = \bar{x}m_b^2$ in the penguin diagram Fig.1.e has well defined meaning and leaves no ambiguity as to the value of k^2, which has usually been treated as a free phenomenological parameter in the estimations of the strong phase generated though the BSS mechanism [45]. So that CP asymmetries could be predicted soundly.

V. B MESON EXCLUSIVE DECAYS TO CHARMONIUM IN QCD FACTORIZATION

by K.T. Chao

Exclusive B-meson decays into charmonium are interesting, since these decays e.g. $B \to J/\psi K$ are regarded as the golden channels for CP violation studies, and these decays provide useful information towards the understanding of perturbative and nonperturbative QCD. These decays are color-suppressed decays, and involve two heavy quark
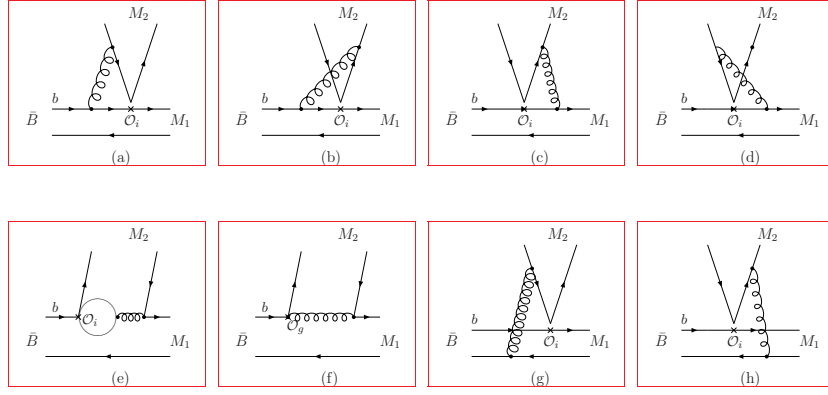
Figure 1:

FIG. 1: Feynman diagrams contributed to non-factorizable contributions.

energy scales, m_b and m_c, therefore are more subtle in theoretical studies. While experimentally CLEO, BaBar and Belle Collaborations have provided many measurements on the B meson exclusive decays to charmonium [46] such as $J/\psi, \psi', \eta_c, \eta_c', \chi_{c0}$, and χ_{c1}, theoretical studies for those decays are still limited.

In the QCD factorization approach [35, 44], it is argued that because the size of charmonium is small $(\sim 1/\alpha_s m_\psi)$ and its overlap with the (B, K) system is negligible in the heavy charm quark limit, QCD factorization method can be used for $B \to J/\psi K$ and other charmonium states, where charmonium is described by the color-singlet $c\bar{c}$ pair. Indeed, explicit calculations for $B \to J/\psi K$ within the QCD factorization approach [47] show that the nonfactorizable vertex contribution is infrared safe and the spectator contribution is perturbatively calculable at twist-2 order, but the theoretical branching ratio is much smaller than the experimental data. Aside from the $B \to J/\psi K$ decay [47], here we present some additional results for B exclusive decays to other S-wave, P-wave, and D-wave charmonium states [48, 49, 50, 51] in QCD factorization. Problem of infrared divergences in B decays to P-wave charmonium $(B \to \chi_{c0} K, B \to \chi_{c2} K,$ and $B \to h_c K)$ and problem of S-D mixing in B decay to D-wave charmonium $(B \to \psi(3770)K)$ are emphasized.

A. $B \to \eta_c K$ decay

Studies of B meson decays into pseudoscalar charmonium states $B \to \eta_c K$ and $B \to \eta_c(2S)K$ [48] in QCD factorization show that the nonfactorizable corrections to naive factorization are infrared safe at leading-twist order. The spectator interactions arising from the kaon twist-3 effects are formally power-suppressed but chirally and logarithmically enhanced. An important improvement by including the $\mathcal{O}(\alpha_s)$ corrections is the large cancellation of the renormalization scale μ dependence of the decay amplitude. However, the calculated decay rates are smaller than experimental data by a factor of 8-10. On the other hand, it is found that for B meson decays to $J/\psi, \psi', \eta_c$ and η_c', the predicted relative decay rates of these four states are approximately compatible with experimental data. For instance, the theoretical value

$$\frac{\text{Br}(B^0 \to \eta_c K^0)}{\text{Br}(B^0 \to J/\psi K^0)_{Th.}} \approx \left(\frac{f_{\eta_c}}{f_{J/\psi}}\right)^2 \cdot \left[\frac{F_1(m_{\eta_c}^2)}{F_1(m_{J/\psi}^2)}\right]^2 \cdot \left[\frac{m_B^2 - m_{\eta_c}^2}{m_B^2 - m_{J/\psi}^2}\right]^3 \approx 1.1 \times \left(\frac{f_{\eta_c}}{f_{J/\psi}}\right)^2 \approx 0.83 - 1.2, \quad (5)$$

is compatible with the experimental value

$$\frac{\text{Br}(B^0 \to \eta_c K^0)}{\text{Br}(B^0 \to J/\psi K^0)_{Ex.}} \approx 1.0. \quad (6)$$

Here F_1 is the $B \to K$ form factor, and the ratio of the squared decay constants $\left(\frac{f_{\eta_c}}{f_{J/\psi}}\right)^2$ ranges from 0.75 to 1.1 in various models.

B. $B \to \chi_{c0}K$, $B \to \chi_{c2}K$, $B \to h_cK$ decays

Studies of B meson decays to P-wave charmonium states $B \to \chi_{c0}K$, $B \to \chi_{c2}K$, and $B \to h_cK$ [49, 50] challenge the applicability of QCD factorization for charmonium because of the appearance of non-vanishing infrared divergences in these decays.

For these decays, there are no leading order contributions from the $V-A$ currents in the effective 4-quark lagrangian, because these P-wave charmonium states do not couple to the $V-A$ currents. According to [35, 44] in QCD factorization all nonfactorizable corrections are due to the vertex corrections and spectator corrections, and other corrections are factorized into the physical form factors and meson wave functions. Taking nonfactorizable corrections into account, the decay amplitude for $B \to \chi_{cJ}K(J = 0, 2)$ can be written as

$$i\mathcal{M} = \frac{G_F}{\sqrt{2}}\Big[V_{cb}V_{cs}^*C_1 - V_{tb}V_{ts}^*(C_4 + C_6)\Big] \times A, \tag{7}$$

where C_i are the well known Wilson coefficients and A is given by

$$A = \frac{i6\mathcal{R}_1^{'}(0)}{\sqrt{\pi M}} \cdot \frac{\alpha_s}{4\pi}\frac{C_F}{N_c}\Big(F_1 \cdot f_I + \frac{4\pi^2 f_K f_B}{N_c} \cdot f_{II}\Big). \tag{8}$$

Here $R_1^{'}(0)$ is the derivative of the radial wave function at the origin of P-wave charmonia , N_c is the number of colors, $C_F = (N_c^2 - 1)/(2N_c)$. The function f_I is calculated from vertex corrections, and f_{II} is calculated from spectator corrections. For $^3P_0$ charmonium state χ_{c0}, the infrared divergence from the vertex correction is found to be

$$f_I = \frac{8m_b z(1 - z + \ln z)}{(1 - z)\sqrt{3z}} \ln\big(\frac{\lambda^2}{m_b^2}\big) + finite\ terms, \tag{9}$$

where $z = M^2/m_B^2 \approx 4m_c^2/m_b^2$, and λ is the gluon mass introduced to regularize the infrared divergences in the vertex corrections. Similar infrared divergences can also be found for the $B \to \chi_{c2}K$ and $B \to h_cK$ decays. The physical reason for the nonappearance of infrared divergences for the S-wave charmonia J/ψ and η_c is based on the fact that the virtual soft gluon in the vertex correction sees a zero total color-charge of the charm quark plus the anti-charm quark; whereas in the P-wave case the virtual soft gluon probes not the color charge of the charm quark and the anti-charm quark but the product of the color charge and the momentum of the charm quark and anti-charm quark, and since the total contribution from the charm quark plus anti-charm quark is nonzero (the relative momentum between the charm quark and anti-charm quark is essential) this leads to the appearance of infrared divergences for the P-wave charmonia. Experimentally, $B \to \chi_{c0}K$ has a large rate, comparable to $B \to \psi'K$ or $B \to \chi_{c1}K$ which are factorizable with the leading order $V-A$ currents. This is certainly very puzzling in the QCD factorization approach.

It is well known that the infrared divergences in the *inclusive* production and decay processes of a color-singlet P-wave $c\bar{c}$ state can be removed by including contributions from the higher Fock states with color-octet $c\bar{c}$ pair (say in S-wave) and soft gluon within the NRQCD factorization framework [52]. However, the color-octet $c\bar{c}$ pair with dynamical soft gluon can make contributions to the multi-body but not two-body *exclusive* decays. As a result, the infrared divergences encountered in *exclusive* processes involving charmonia may raise a new question to the QCD factorization and NRQCD factorization in B exclusive decays. Further studies are needed to seek the solution to remove the infrared divergences.

As for the leading order factorizable $B \to \chi_{c1}K$ decay, the calculated rate is again too small [49], as in other leading order factorizable cases (e.g. $B \to J/\psi K$ and $B \to \eta_cK$).

C. $B \to \psi(3770)K$ decay

The $\psi(3770)$ is the $J^{PC} = 1^{--}$ 1D-wave charmonium with a small admixture of 2S component, and the mixing angle is estimated to be about $\theta = -12°$ [53]. Recently, Belle [54] has observed $\psi(3770)$ in the B meson decay $B^+ \to \psi(3770)K^+$ with a large branching ratio of $(0.48 \pm 0.11 \pm 0.07) \times 10^{-3}$, which is comparable to $\mathcal{B}(B^+ \to \psi'(3686)K^+) = (0.66 \pm 0.06) \times 10^{-3}$ [46].

The $B \to \psi(3770)K$ decay rate is calculated with QCD radiative corrections in QCD factorization [51] and it is found that the 1D contribution is much smaller than 2S though both 1D and 2S components in the $\psi(3770)$ wave function can couple to the leading order vector current and that a large mixing angle $\theta = -26°$ is required to fit the observed ratio $\frac{Br(B \to \psi'' K)}{Br(B \to \psi' K)}$. The required large mixing angle is in contradiction with all other estimates in charmonium phenomenology [53]. Moreover, the calculated $B \to \psi(3770)K$ decay rate is still smaller than data by a factor of 8

even with the large mixing angle included. In contrast, in the B inclusive decays the D-wave charmonium states are expected to have large rates due to the color-octet mechanism regardless of the S-D mixing angle [55].

In summary, there are a number of problems in the description of B meson exclusive decays to charmonium states in the QCD factorization approach. The decay rates for the S-wave states such as J/ψ and η_c are too small to accommodate the experimental data. For the exclusive decays to P-wave charmonia $B \to \chi_{cJ}K$, it is found that (except χ_{c1}) there are infrared divergences arising from nonfactorizable vertex corrections as well as logarithmic end-point singularities arising from nonfactorizable spectator interactions at leading-twist order. The infrared divergences due to vertex corrections will explicitly break down QCD factorization within the color-singlet model for charmonia. Unlike in the inclusive decays where the higher Fock states with color-octet $c\bar{c}$ pair and soft gluon can make contributions to remove the infrared divergences, their contributions can not be accommodated in the exclusive two body decays. As a result, the infrared divergences encountered in exclusive processes involving charmonia may raise a new question to the QCD factorization and NRQCD factorization in B exclusive decays. For the D-wave state the observed $B \to \psi(3770)K$ decay rate would require a much larger S-D mixing angle than conventional expected in charmonium phenomenology. All these problems are likely to be related to the limitation of the color-singlet $c\bar{c}$ pair description for charmonium in QCD factorization. But how to solve these problems still remains an open question. New theoretical considerations or ingredients in QCD are needed.

VI. PERTURBATIVE QCD APPROACH BASED ON k_T FACTORIZATION

by C.D. Lü, X.Q. Yu

Understanding non-leptonic B meson decays is crucial for testing the standard model, and also for uncovering the trace of new physics. The simplest case is two-body non-leptonic B meson decays, for which Bauer, Stech and Wirbel (BSW) proposed the naive factorization assumption (FA) in their pioneering work [39]. The technique to analyze hard exclusive hadronic scattering was developed by Brodsky and Lepage [24] based on collinear factorization theorem in perturbative QCD (PQCD). A modified framework based on k_T factorization theorem was then given in [56, 57], and extended to exclusive B meson decays in [58, 59, 60, 61]. The infrared finiteness and gauge invariance of k_T factorization theorem was shown explicitly in [62]. Using this so-called PQCD approach, we have investigated dynamics of non-leptonic B meson decays [34, 63, 64]. Our observations are summarized as follows:

1. FA holds approximately for many charmless B meson decays, as our computation shows that non-factorizable contributions are always negligible due to the cancellation between a pair of non-factorizable diagrams.

2. Penguin amplitudes are enhanced, as the PQCD formalism includes dynamics from the region, where the energy scale μ runs to $\sqrt{\bar{\Lambda} m_b} < m_b/2$, $\bar{\Lambda} \equiv m_B - m_b$ being the B meson and b quark mass difference.

3. Annihilation diagrams contribute to large short-distance strong phases through the $(S + P)(S - P)$ penguin operators.

4. The sign and magnitude of CP asymmetries in two-body non-leptonic B meson decays can be calculated, and we have predicted relatively large CP asymmetries in the $B \to K^{(*)}\pi$ [34, 65] and $\pi\pi$ modes [63, 64, 66].

A. Formalism of PQCD Approach

To develop the PQCD formalism for charmed B meson decays, we investigated the $B \to D^{(*)}$ transition form factors in the large recoil region of the $D^{(*)}$ meson [68]. The $B \to D^{(*)}$ transition is more complicated than the $B \to \pi$ one, because it involves three scales: the B meson mass m_B, the $D^{(*)}$ meson mass $m_{D^{(*)}}$, and the heavy meson and heavy quark mass difference, $\bar{\Lambda} = m_B - m_b \sim m_{D^{(*)}} - m_c$ of order of the QCD scale Λ_{QCD}, $m_{D^{(*)}}$ (m_c) being the $D^{(*)}$ meson (c quark) mass. We have postulated the hierarchy of the three scales,

$$m_B \gg m_{D^{(*)}} \gg \bar{\Lambda} \,, \tag{10}$$

which allows a consistent power expansion in $m_{D^{(*)}}/m_B$ and in $\bar{\Lambda}/m_{D^{(*)}}$.

In the $B \to D^{(*)}$ transition , the initial state is approximated by the $b\bar{d}$ component. The b quark decays into a c quark and a virtual W boson, which carries the momentum q. Since the constituents are roughly on the mass shell, we have the invariant masses $k_i^2 \sim O(\bar{\Lambda}^2)$, $i = 1$ and 2, where k_1 (k_2) is the momentum of the spectator $\bar{d}$ quark in the B ($D^{(*)}$) meson. The lowest-order diagrams contributing to the $B \to D^{(*)}$ form factors contain a hard gluon

Quantities	$C_D = 0.6$	$C_D = 0.8$	$C_D = 1.0$	Data
$B(\bar{B}^0 \to D^+\pi^-)$	2.37	2.74	3.13	3.0 ± 0.4
$B(\bar{B}^0 \to D^0\pi^0)$	0.26	0.25	0.24	0.29 ± 0.05
$B(B^- \to D^0\pi^-)$	4.96	5.43	5.91	5.3 ± 0.5
Decay Modes	$C_{D*} = 0.5$	$C_{D*} = 0.7$	$C_{D*} = 0.9$	Data
$B(\bar{B}^0 \to D^{*+}\pi^-)$	2.16	2.51	2.88	2.76 ± 0.21
$B(\bar{B}^0 \to D^{*0}\pi^0)$	0.29	0.28	0.27	0.25 ± 0.07
$B(B^- \to D^{*0}\pi^-)$	4.79	5.26	5.75	4.60 ± 0.40

TABLE I: Predicted $B \to D^{(*)}\pi$ decay branching ratios in units of 10^{-3}.

exchange between the b or c quark and the $\bar{d}$ quark. The $\bar{d}$ quark undergoes scattering in order to catch up with the c quark, forming a $D^{(*)}$ meson. The exchanged gluon is off-shell by

$$(k_1 - k_2)^2 \sim -\frac{m_B}{m_{D^{(*)}}}\bar{\Lambda}^2 \,, \tag{11}$$

which has been identified as the characteristic scale of the hard kernels. Under Eq. (10), we have $m_B/m_D^{(*)} \gg 1$, and the hard kernels are calculable in perturbation theory. It has been found that the applicability of PQCD to the $B \to D^{(*)}$ transition at large recoil is marginal for the physical masses m_B and $m_{D^{(*)}}$ [68].

The form factors are then expressed as the convolution of the hard kernels H with the B and $D^{(*)}$ meson wave functions in k_T factorization theorem,

$$F^{BD^{(*)}}(q^2) = \int dx_1 dx_2 d^2b_1 d^2b_2 \phi_B(x_1, b_1) H(x_1, x_2, b_1, b_2) \phi_{D^{(*)}}(x_2, b_2) \,. \tag{12}$$

The $D^{(*)}$ meson wave function contains a Sudakov factor arising from k_T resummation, which sums the large double logarithms $\alpha_s \ln^2(m_B b_2)$ to all orders. The B meson wave function also contains such a Sudakov factor, whose effect is negligible because a B meson is dominated by soft dynamics. The hard kernels involve a Sudakov factor from threshold resummation, which sums the large double logarithm $\alpha_s \ln^2 x_1$ or $\alpha_s \ln^2 x_2$ to all orders. This factor modifies the end-point behavior of the B and $D^{(*)}$ meson wave functions effectively, rendering them diminish faster in the small $x_{1,2}$ region.

It has been pointed out that if evaluating the hard part with a hard gluon in collinear factorization theorem, end-point singularities appear [69]. These singularities imply the breakdown of collinear factorization, and k_T factorization becomes more appropriate. Once the parton transverse momenta k_T are taken into account, a dynamical effect, the so-called Sudakov suppression, favors the configuration in which k_T is not small [57]. The end-point singularities then disappear. The Sudakov factor suppresses the large b region, where the quark and anti-quark are separated by a large transverse distance and the color shielding is not so effective. It also suppresses the $x \sim 1$ region, where a quark carries all of the meson momentum, and intends to emit real gluons in hard scattering.

B. $B \to D^{(*)}\pi$ IN PQCD

In this subsection we take the $B \to D\pi$ decays as an example of the PQCD analysis. In the PQCD framework based on k_T factorization theorem, an amplitude is expressed as the convolution of hard b quark decay kernels with meson wave functions in both the longitudinal momentum fractions and the transverse momenta of partons. Our PQCD formulas are derived up to leading-order in α_s, to leading power in m_D/m_B and in $\bar{\Lambda}/m_D$, and to leading double-logarithm resummations.

The predicted branching ratios in Table I are in agreement with the averaged experimental data [71, 72, 73]. We extract the effective parameters a_1 and a_2 from our PQCD calculations. That is, our a_1 and a_2 do not only contain the non-factorizable amplitudes as in generalized FA, but the small annihilation amplitudes, which was first discussed in [74]. We obtain the ratio $|a_2/a_1| \sim 0.43$ with 10% uncertainty and the phase of a_2 relative to a_1 about $Arg(a_2/a_1) \sim -42°$. Note that the experimental data do not fix the sign of the relative phases. The PQCD calculation indicates that $Arg(a_2/a_1)$ should be located in the fourth quadrant. It is evident that the short-distance strong phase from the color-suppressed non-factorizable amplitude is already sufficient to account for the isospin triangle formed by the $B \to D\pi$ modes. From the viewpoint of PQCD, this strong phase is of short distance, and produced from the

TABLE II: PQCD predictions with one angle mixing formalism (I) and two angle mixing formalism (II) and experimental data (in units of 10^{-4}) of the $B^0 \to \bar{D}^{(*)0}\eta^{(\prime)}$ branching ratios.

Decay mode	PQCD (I)	PQCD (II)	Belle	BaBar	PDG
$B^0 \to \bar{D}^0\eta^{\prime}$	$1.7 \sim 2.3$	$2.2 \sim 2.6$	-	-	< 9.4
$B^0 \to \bar{D}^0\eta$	$2.4 \sim 3.0$	$2.6 \sim 3.2$	$1.4^{+0.6}_{-0.5}$	2.41 ± 0.50	
$B^0 \to \bar{D}^{*0}\eta^{\prime}$	$2.0 \sim 2.7$	$2.6 \sim 3.2$		-	< 14
$B^0 \to \bar{D}^{*0}\eta$	$2.8 \sim 3.5$	$3.1 \sim 3.8$	$2.0^{+1.0}_{-0.9}$	-	

non-pinched singularity of the hard kernel. Certainly, under the current experimental and theoretical uncertainties, there is still room for long-distance phases from final-state interaction.

The PQCD predictions for the $B \to D^*\pi$ decay branching ratios in Table I are also consistent with the data [46]. Since m_{D^*} and ϕ_{D^*} are slightly different from m_D and ϕ_D, respectively, the results are close to those for $B \to D\pi$. Similarly, the ratio $|a_2/a_1|$ and the relative phase $Arg(a_2/a_1)$ are also close to those associated with the $B \to D\pi$ decays. We obtain the ratio $|a_2/a_1| \sim 0.47$ with 10% uncertainty and the relative phase about $Arg(a_2/a_1) \sim -41°$. The $B \to D^{(*)}\eta^{(\prime)}$ decays are similar to the $B \to D^{(*)}\pi^0$ decays except the η -$\eta^{\prime}$ mixing part [75]. The predicted branching ratios for these decays are shown in table II, with uncertainties only from the η -$\eta^{\prime}$ mixing parameter θ.

C. $B_s \to \pi^+\pi^-$ decay IN PQCD

We also calculate the decay rate and CP asymmetry of the $B_s \to \pi^+\pi^-$ decay in perturbative QCD approach with Sudakov resummation. Since none of the quarks in final states is the same as those of the initial B_s meson, this decay can occur only via annihilation diagrams in the standard model. Besides the current-current operators, the contributions from the QCD and electroweak penguin operators are also taken into account. We find that (a) the branching ratio is about 4×10^{-7}; (b) the penguin diagrams dominate the total contribution; and (c) the direct CP asymmetry is small in size: no more than 3%; but the mixing-induced CP asymmetry can be as large as ten percent testable in the near future LHC-b experiments and BTeV experiments at Fermilab. This small branching ratio, predicted in the SM, make it sensitive to possible new physics contribution.

VII. B MESON DECAYS AND THE CKM ANGLES α, β AND γ

by Z.J. Xiao

For discussion on isospin and SU(3) relations in charmless B decays and possible new type of electroweak penguin or SU(3) breaking effects of strong phases, as well as large direct CP violation in charmless B decays, it is suggested to see ref. [22]. As a summary, we observe that in the case of SU(3) limits and also the case with SU(3) breaking only in amplitudes, the fitting results lead to an unexpected large ratio between two isospin amplitudes $a^c_{3/2}/a^u_{3/2}$, which is about an order of magnitude larger than the SM prediction. The results are found to be insensitive to the weak phase γ. By including SU(3) breaking effects on the strong phases, one is able to obtain a consistent fit to the current data within the SM, which implies that the SU(3) breaking effect on strong phases may play an important role in understanding the observed charmless hadronic B decay modes $B \to \pi\pi$ and πK. It is possible to test those breaking effects in the near future from more precise measurements of direct CP violation in B factories.

The direct CP violation with SU(3) symmetry breaking of strong phases can be as large as follows:

$$A_{CP}(\pi^+\pi^-) \simeq 0.5, \qquad A_{CP}(\pi^0\pi^0) \simeq 0.2, \qquad A_{CP}(\pi^+K^-) \simeq -0.1,$$
$$A_{CP}(\pi^0\bar{K}^0) \simeq -0.1, \qquad A_{CP}(\pi^0K^-) \simeq -0.0, \qquad A_{CP}(\pi^-\bar{K}^0) \simeq 0.1.$$

For CP violation in $B \to \phi K_S$ and possible new physics models, see ref. [23] and references therein.

One of the most important tasks in B factory experiments is to measure the CKM angles α, β and γ [76], the three inner angles of the unitarity triangle defined in the complex $(\bar{\rho}, \bar{\eta})$ plane. At present, $B \to \pi\pi, K\pi$ and many other interesting decays have been measured with good precision in B factory experiments. These modes together with the CP asymmetries of $B \to J/\Psi K_S, \rho\rho, \rho^{\pm}\pi^{\mp}, etc$ will allow us to determine the CKM angles α, β and γ. The isospin symmetry of strong interaction and $SU(3)$ flavor symmetry are frequently used to deal with "unknown" hadronic matrix elements, and in some cases also plausible dynamical assumptions have to be made.

From a global fit done by CKM fitter group [77] the currently allowed ranges of these three angles at 95% C.L. are

$$78° \leq \alpha \leq 122°, \quad 21° \leq \beta \leq 27°, \quad 38° \leq \gamma \leq 80° \tag{13}$$

The angle β has already been well measured using the time-dependent CP asymmetries in the $B_d^0 \to J/\Psi K_S$ and related decays. The world average as given by Heavy Flavor Averaging Group [78] is

$$\begin{aligned} \sin 2\beta &= 0.739 \pm 0.048, \quad all \ \ charmonium, \\ &= 0.692 \pm 0.045, \quad all \ \ modes \end{aligned} \tag{14}$$

which leads to the bounds on the angle β:

$$\begin{aligned} \beta &= \left(23.8^{+2.2}_{-2.0}\right)° \bigvee \left(66.2^{-2.2}_{+2.0}\right)°, \quad all \ \ charmonium, \\ &= \left(21.9^{+1.9}_{-1.7}\right)° \bigvee \left(68.1^{-1.9}_{+1.7}\right)°, \quad all \ \ modes \end{aligned} \tag{15}$$

The twofold ambiguity has been resolved by BaBar's new measurement of $\cos 2\beta$: a negative $\cos 2\beta$ is excluded at 89%C.L. [79]. We therefore believe that $\beta = 24° \pm 2°$ is a reliable measurement and can be used as an experimental input in the effort to determine other two angles α and γ.

In the SM, the penguin-dominated $B \to \phi K_s$ decay should measure the same $\sin 2\beta$ as $B \to J/\Psi K_S$ decay. However, a 3.5σ deviation from the SM prediction is observed by Belle collaboration. The most recent measurements give [80]

$$\sin 2\beta|_{\phi K_s} = \begin{cases} -0.96 \pm 0.50^{+0.09}_{-0.11} & (Belle) \\ +0.47 \pm 0.34^{+0.08}_{-0.06} & (BaBar) \end{cases} \tag{16}$$

Although it is too early to draw any conclusion from above primary measurements, many works have been done to interpret such deviation as a hint of new physics contributions to the quark level $b \to ss\bar{s}$ decays.

The CKM angle α can be extracted through the CP violation in the tree-dominated $B \to \pi\pi, \rho\pi$ and $B \to \rho\rho$ decays, The $B \to \pi^+\pi^-$ decay currently plays most important rule in constraining α. The isospin symmetry can be used to eliminate the penguin pollution, but a model-independent isospin analysis of $B \to \pi\pi$ decays requires the knowledge of the three amplitudes A^{+-}, A^{00}, A^{+0} and their charge conjugates $\bar{A}^{ij}$, At present, the only missing pieces are A^{00} and $\bar{A}^{00}$. It is possible that a full isospin analysis will be done, and α extracted cleanly, by the summer of 2005.

Theoretically, the CP asymmetry parameter $C_{\pi\pi}$ and $S_{\pi\pi}$ depend on α, the strong phase δ and the ratio $r = |P/T|$:

$$S_{\pi\pi} = \frac{\sin 2\alpha - 2r\cos\delta\sin\alpha}{1 + r^2 - 2r\cos\delta\cos\alpha}, \quad C_{\pi\pi}\frac{2r\sin\delta\sin\alpha}{1 + r^2 - 2r\cos\delta\cos\alpha}. \tag{17}$$

In ref. [81], taking $S_{\pi\pi} = -0.49 \pm 0.27$ and $C_{\pi\pi} = 0.51 \pm 0.19$ as experimental input, we found the allowed ranges: (a) $76° \leq \alpha \leq 135°$ for $r = 0.3 \pm 0.1$; and (b) $117° \leq \alpha \leq 135°$ and $-160° \leq \delta \leq -132°$ if we also take the measured branching ratios of $B \to \pi^+\pi^-, \pi^+\pi^0$ and $K^0\pi^+$ decays into account.

In ref. [82], by setting f_K/f_π as the $SU(3)$ flavor symmetry breaking factor, and using the measured branching ratios of $B \to \pi^+\pi^-$ and $K\pi$ decays and the new average $S_{\pi\pi} = -0.74 \pm 0.16$ and $C_{\pi\pi} = -0.46 \pm 0.13$ [78] as input, Gronau found the allowed range, $\alpha = (104 \pm 18)°$ [82]. From currently available measurements of $B \to \rho\rho$ decays, the range $19° \leq \alpha \leq 71°$ is excluded at 90%C.L. [82]. From the measured time-dependent CP asymmetries in $B \to \rho^\pm\pi^\mp$, one found that the allowed range is $\alpha = 93° \pm 17°$ ($\alpha = 102° \pm 11°$) if only the BaBar (Belle) measurements are taking into account.

By using the SM relation $\alpha + \beta + \gamma = 180°$, the above bounds on α can be translated into bounds on γ for fixed value of $\beta = 24° \pm 2°$. Of course, many strategies to determine angle γ directly from B meson decays have been proposed recently [76, 83, 84, 85], for example,

- Isospin symmetry plus the $B \to K\pi$ observables allow one to constrain the angle γ. In ref. [84], for example, the region $\gamma \sim 90°$ was excluded by considering a "mixed" system of $B^+ \to \pi^+ K^0$ and $B_d^0 \to \pi^0 K^0$ decays.

- U-spin symmetry plus the measurements of $B_d \to \pi\pi$ and $B_s \to KK$, , $B_d \to J/\Psi K_S$ and $B_s \to D_s^+ D_s^-$, or $B_s \to D_s^{(*)\pm} K^\mp$ and $B_d \to D^{(*)\pm}\pi^\mp$, etc.

- to determine α through the measurements of the branching ratios of six $B \to DK$ decays.

VIII.   INSTANTANEOUS BETHE-SALPETER EQUATION AND ITS EXACT SOLUTIONS

by C.H. Chang

We propose an approach to solve a Bethe-Salpeter (BS) equation exactly without approximation if its kernel exactly is instantaneous, and take positronium as an example to illustrate the general features of the approach and the obtained solutions under it [86].

The key step for the approach is to start with the BS equation to derive out a set of coupled and well-determined integration equations for the components of the BS wave function without any approximation as a linear eigenvalue problem. The coupled equations may be solved analytically and/or numerically depending on the specific kernel. If there is no analytic solution, in principle, the coupled equations can be also solved numerically under a controlled (requested) accuracy.

To solve the coupled and well-determined equations for positronium, a numerical method, which is to expand the equations (the eigenfunctions and the kernel) for the BS wave function components in terms of the relevant Schrödinger equation solutions with suitable truncation and then alternatively to solve the truncated matrix equation by making it diagonal, is applied, and accurate enough solutions for requests (the accuracy depends on the truncation mainly) are obtained. The exact solutions present precise and substantial corrections to those of the corresponding Schrödinger equation, which start from the order $\mathcal{O}(v^1)$ (v is the relative velocity) for eigenfunctions, the order $\mathcal{O}(v^2)$ for eigenvalues and the mixing between S (P) and D (F) components in $J^{PC} = 1^{--}$ ($J^{PC} = 2^{++}$) states etc. From the lessons, we point out that such corrections are important in the effective theories such as NRQCD and NRQED when we consider the relativistic corrections $\mathcal{O}(v) \sim \mathcal{O}(\frac{1}{M})$. Namely **one cannot 'forget' the same order corrections involved in the solutions of the bound states accordingly.**

Moreover, we also point out that there is a questionable step in the classical derivations for an instantaneous BS equation, if one is pursuing the exact solutions without approximation. Namely it has confused the differences between the instantaneous BS equation and Breit equation (more precise illustration in [87]).

IX.   NEW PHYSICS EFFECTS IN B DECAYS

by X.G. He and C.S. Huang and Z.J. Xiao

A.   Time Dependent CP Violation in $B \to K^*\gamma$ in SUSY

In this work we study time dependent CP asymmetries in $B \to K^*\gamma \to \pi K_S \gamma$ and $B \to \phi K_S$ using constraints from $B \to X_S\gamma$ in SUSY model with large Left-Right (or Right-Left) squark mixing induced gluonic dipole interaction. Since this work was reported at the meeting, new results on B decays have been reported [88, 89, 90, 91], in particular the direct CP asymmetry A_{CP} in $\bar{B}^0 \to K^-\pi^+$ of [88] -0.114 ± 0.020 measured by Babar and Belle, we include the constraint from $A_{CP}(K^-\pi^+)$ in our study [92].

In the SM, the Hamiltonian for the B decays to be considered is well known which is of the form [93],

$$H = \frac{G_F}{\sqrt{2}}[V_{ub}V_{us}^*(c_1 O_1 + c_2 O_2) - \sum_{i=3}^{12} V_{jb}V_{js}^* c_i^j O_i], \tag{18}$$

where V_{ij} are the CKM matrix elements. c_i are the Wilson coefficients for the operators O_i which have been evaluated in different schemes. Values from NDR scheme will be used [93]. We will not display the full sets of O_i and c_i here, but only give the definitions of the gluonic and photonic dipole operators O_{11} and O_{12} for the convenience of later discussions. They are given by

$$
\begin{aligned}
O_{11} &= \frac{g_s}{8\pi^2}\bar{s}\sigma_{\mu\nu}G_a^{\mu\nu}T^a[m_b(1+\gamma_5)+m_s(1-\gamma_5)]b, \\
O_{12} &= \frac{e}{8\pi^2}\bar{s}\sigma_{\mu\nu}F^{\mu\nu}[m_b(1+\gamma_5)+m_s(1-\gamma_5)]b,
\end{aligned}
\tag{19}
$$

where T^a is the color SU(3) generator normalized to $Tr(T^a T^b) = \delta^{ab}/2$. $G_{\mu\nu}$ and $F_{\mu\nu}$ are the gluon and photon field strengths. In the SM [93, 94, 95] $c_{11} = -0.151$ and $c_{12} = -0.318$.

When going beyond the SM, there are modifications. In SUSY models, exchanges of gluino and squark with Left-Right squark mixing can generate a large contribution to $c_{11,12}$ at one loop level [96, 97] since their interactions are

strong couplings in strength and also enhanced by a factor of the ratio of gluino mass to the b quark mass [98]. We will concentrate on the effects of this interaction. In general exchange of squarks and gluinos can generate non-zero $c_{11,12}$ for dipole operators with $1 + \gamma_5$, as well as with non-zero $c'_{11,12}$ for dipole operators with $1 - \gamma_5$.

The Wilson coefficient $c_{11,12}^{susy}$ for SUSY contribution obtained in the mass insertion approximation is given by, for the case with $1 + \gamma_5$ [97],

$$c_{11}^{susy}(m_{\tilde{g}}) = \frac{\sqrt{2}\pi\alpha_s(m_{\tilde{g}})}{G_F m_{\tilde{g}}^2} \frac{\delta_{LR}^{bs}}{V_{tb}V_{ts}^*} \frac{m_{\tilde{g}}}{m_b} G_0(x_{gq}),$$

$$c_{12}^{susy}(m_{\tilde{g}}) = \frac{\sqrt{2}\pi\alpha_s(m_{\tilde{g}})}{G_F m_{\tilde{g}}^2} \frac{\delta_{LR}^{bs}}{V_{tb}V_{ts}^*} \frac{m_{\tilde{g}}}{m_b} F_0(x_{gq}),$$

$$G_0(x) = \frac{x[22 - 20x - 2x^2 + (16x - x^2 + 9)\ln(x)]}{3(1-x)^4},$$

$$F_0(x) = -\frac{4x[1 + 4x - 5x^2 + (4x + 2x^2)\ln(x)]}{9(1-x)^4}, \tag{20}$$

where δ_{LR}^{bs} parameterizes the mixing of left and right squarks, $x_{gq} = m_{\tilde{g}}^2/m_{\tilde{q}}^2$ is the ratio of gluino mass $m_{\tilde{g}}$ and squark mass $m_{\tilde{q}}$. The Wilson coefficients $c_{11,12}^{'susy}$ for the case with $1 - \gamma_5$ can be easily obtained by replacing the Left-Right mixing parameter δ_{LR}^{bs} by the Right-Left mixing parameter δ_{RL}^{bs}.

From the expressions in eq.(20), one can see that the SUSY contributions are proportional to $m_{\tilde{g}}$. If $m_{\tilde{g}}$ is of order a few hundred GeV, there is an enhancement factor of $(m_{\tilde{g}}/m_b)(m_W^2/m_{\tilde{g}}^2)$ for the SUSY dipole interactions. In this case even a small $\delta_{LR,RL}^{bs}$, which can easily satisfy constraints from $B - \bar{B}$ mixing and other data, can have a large effect on rare B decays.

We first consider constraint on the SUSY parameters $\delta_{LR,RL}^{bs}$ from $B \to X_S\gamma$. The branching ratio of this process has been measured to a good precision with [89] $(3.54^{+0.30}_{-0.28}) \times 10^{-4}$. Although experimentally CP asymmetry in $B \to X_S\gamma$ has not been established, there are constraints from experiments with [90] 0.005 ± 0.036. We follow ref.[95] to carry out the evaluation of the branching ratio and CP asymmetry and use them to constrain the parameters.

In Figure 2 we show the allowed ranges for the absolute value of δ_{LR}^{bs} and its phase τ for $m_{\tilde{g}} = 300$GeV and $m_{\tilde{q}}$ in the range $100 \sim 1000$GeV at the one σ level. We find that the constraints from $BR(B \to X_S\gamma)$ and $A_{CP}(B \to X_S\gamma)$ are similar. Using the allowed parameters, one can obtain the allowed c_{11} through eq.20 and to study implications for other rare B decays.
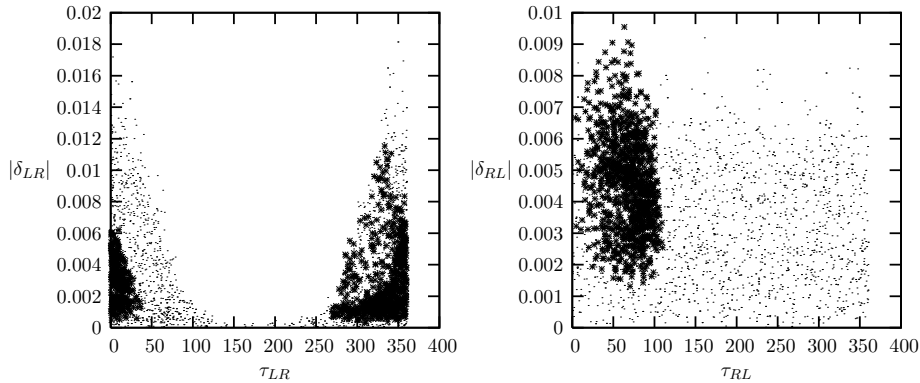

FIG. 2: The one σ allowed ranges for the SUSY parameters $|\delta_{LR,RL}^{bs}|$ and the phase τ taking $m_{\tilde{g}} = 300$GeV and $m_{\tilde{q}}$ in the range $100 \sim 1000$ GeV. The light-dark doted areas are the allowed parameter spaces from $BR(B \to X_s\gamma)$ and $A_{CP}(B \to X_S\gamma)$ constraints. The dark doted areas are allowed ranges by $A_{CP}(\bar{B}^0 \to K^-\pi^+)$ constraint. Figure 1a (on the left) and Figure 1b (on the right) are for the dipole operators with $1 + \gamma_5$ and $1 - \gamma_5$, respectively.

The recently measured CP asymmetry $A_{CP}(\bar{B}^0 \to K^-\pi^+)$ can be reproduced and the SUSY parameters can be further constrained. We follow ref.[99] to calculate the SUSY dipole contribution to $B \to K\pi$. In our numerical analysis we take the CKM parameters to be known, with the standard parametrization $s_{12} = 0.2243$, $s_{23} = 0.0413$, $s_{13} = 0.0037$, $\delta_{13} = 1.05$, which is the central value given by the Particle Data Group [46]. With the SM amplitudes obtained and the default values for the hadronic parameters used in Ref. [99], we obtain the CP asymmetry $A_{CP}(\bar{B}^0 \to$

$K^-\pi^+$) in the SM to be 0.15. This is different in sign with the experimental value. When SUSY dipole interactions are included the experimental value can be reproduced. For example with $m_{\tilde{g}} = m_{\tilde{q}} = 300\text{GeV}$, $\delta_{LR} = 2.62 \times 10^{-3}e^{0.238i}$, $\delta_{RL} = 4.31 \times 10^{-3}e^{1.007i}$ the asymmetry $A_{CP}(\bar{B}^0 \to K^-\pi^+)$ is approximately -0.114. With the same set of SUSY parameters, we have $Br(B \to X_s\gamma) = 3.48 \times 10^{-4}$, $A_{CP}(B \to X_S\gamma) = 0.016$. It is clear that the CP asymmetry $A_{CP}(\bar{B}^0 \to K^-\pi^+)$ can be brought to be in agreement with data at one sigma level when SUSY gluonic dipole interactions are included.

To see how the CP asymmetry provides stringent constraint on the SUSY flavor changing parameters, we show in Figure 2 the parameter space allowed from $A_{CP}(\bar{B}^0 \to K^-\pi^+)$ (the dark doted areas) on top of the allowed ranges by $B \to X_S\gamma$ constraint alone at the one σ level. We see that the CP asymmetry in $\bar{B}^0 \to K^-\pi^+$ considerably reduces the allowed parameter space.

We now study time dependent CP asymmetries in $B \to K^*\gamma \to \pi^0 K_S\gamma$ and $B \to \phi K_S$. There are two CP violating parameters A_f and S_f which can be measured in time dependent decays of B and $\bar{B}$ produced at e^+e^- colliders at the $\Upsilon(4S)$ resonance, $A^{CP}(t) = A_f \cos(\Delta t \Delta m_B) + S_f \sin(\Delta t \Delta m_B)$.

For $\bar{B}^0 \to \bar{K}^{*0}\gamma \to \pi^0 K_S\gamma$ and $B^0 \to K^{*0}\gamma \to \pi^0 K_S\gamma$, we obtain [95, 98, 100]

$$S_{K^*\gamma} = -2\frac{Im[(q_B/p_B)(c_{12}c'_{12})]}{|c_{12}|^2 + |c'_{12}|^2}. \tag{21}$$

To the leading order $A_{K^*\gamma}$ is the same as $A_{CP}(B \to X_S\gamma)$. Note that the hadronic matrix element $< K^*|\bar{s}\sigma^{\mu\nu}(1 \pm \gamma_5)b|B >$ does not appear making the calculation simple and reliable. In order to have a non-zero $S_{K^*\gamma}$ both c_{12} and c'_{12} cannot be zero.

In the SM the asymmetries $A_{K^*\gamma}$ and $S_{K^*\gamma}$ are predicted to be small with [95, 100] $A_{K^*\gamma}^{SM} \approx 0.5\%$, $S_{K^*\gamma}^{SM} \approx 3\%$. With SUSY gluonic dipole interaction, the predictions for these CP asymmetries can be changed dramatically [98]. With the constraints obtained previously, we find that the parameter q_B/p_B is not affected very much compared with the SM calculation. To a good approximation $q_B/p_B = e^{-2i\beta}$.

A Large gluonic dipole interaction also has a big impact on $B \to \phi K_S$ decays [101]. In the SM, $A_{\phi K_s}$ is predicted to be very small and $S_{\phi K_S}$ is predicted to be the same as $S_{J/\psi K_S} = \sin(2\beta)$. With SUSY gluonic dipole contribution, the decay amplitude for $B \to \phi K_S$ will be changed and the predicted value for both $A_{\phi K_S}$ and $S_{\phi K_S}$ can be very different from those in the SM [101]. We again use QCD factorization [99] to evaluate the amplitude.
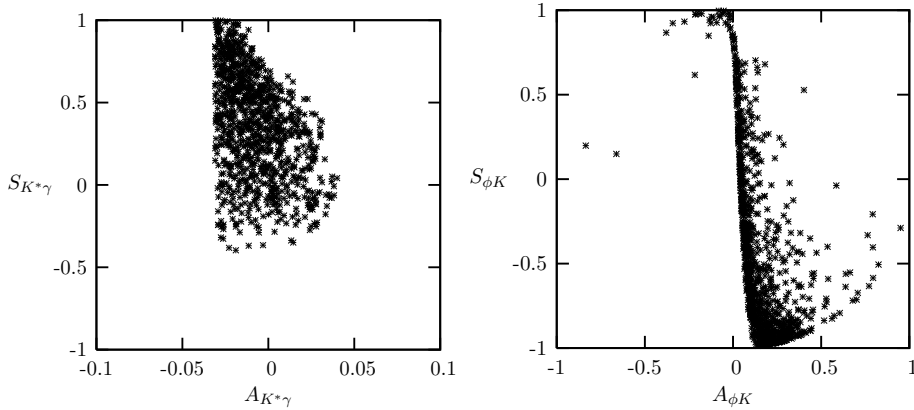

FIG. 3: The allowed time dependent CP asymmetries in $\bar{B}^0 \to K^*\gamma \to K_S\pi^0\gamma$ and $\bar{B}^0 \to \phi K_S$.

The results are shown in Figure 3. The current values of $S_{K^*\gamma}$ and $A_{K^*\gamma}$ from Babar (Belle) are [102]: $0.57 \pm 0.32 \pm 009(-0.00 \pm 0.38)$, $0.25 \pm 0.63 \pm 0.14(-0.79^{+0.63}_{-0.50} \pm 0.09)$, respectively. From Figure 3, we see that the allowed ranges can cover the central values of $S_{K^*\gamma}$ from Babar and Bell, but it is not possible to obtain the central value of $A_{K^*\gamma}$ by Belle. Future improved data can further restrict the parameter space. Both Babar and Belle have also measured [104] $A_{CP}(B^- \to K^{*-}\gamma)$ with ranges $-0.074 \sim 0.049$(Babar) and $-0.015 \pm 0.044 \pm 0.012$(Belle). In the model we are considering, the CP asymmetries $A_{K^*\gamma}$ and $A_{CP}(B^- \to K^{*-}\gamma)$ are the same. The results for the charged B CP asymmetry is consistent with data.

The time dependent asymmetry in $B \to \phi K_S$ is a very good test of CP violation in the SM. Experimental measurements have not converged with the current values of Babar (Belle) given by [91, 103] $0.00\pm0.23\pm0.05(0.08\pm0.22\pm0.09)$, and $0.50\pm0.25^{+0.07}_{-0.04}(0.06\pm0.33\pm0.09)$ for $A_{\phi K_S}$ and $S_{\phi K_S}$, respectively. These values are considerably different than

the value reported by Belle last year of [105] $S_{\phi K_S} = -0.96 \pm 0.50^{+0.09}_{-0.11}$. From Figure 3 we see that the current data of $A_{\phi K_S}$ and $S_{\phi K_S}$ can be easily accommodated by the allowed ranges. We also note that the allowed ranges can cover last year's Belle data. Since the error bars on the data are large, no definitive conclusions can be drawn at present.

B. CP asymmetries in $B \to \phi K_S$ and $B \to \eta' K_S$ in MSSM

The time dependent CP asymmetry in $B \to \phi K_S$

$$
\begin{aligned}
S_{\phi K_S} &= -0.39 \pm 0.41, \quad 2002 \ \text{World} - \text{average} \\
S_{\phi K_S} &= -0.15 \pm 0.33, \quad 2003 \ \text{World} - \text{average}
\end{aligned}
\tag{22}
$$

is especially interesting since it deviates greatly from the SM expectation

$$
S_{\phi K_S} = \sin(2\beta(\phi K_S)) = \sin(2\beta(J/\psi K_S)) + O(\lambda^2)
\tag{23}
$$

where $\lambda \simeq 0.2$ appears in Wolfenstein's parametrization of the CKM matrix. If the error in the average is inflated by the scale factor, the 2003 World-average will be -0.15 ± 0.69, which deviates from the SM only by 1.3σ, i.e., the statistical basis of the effect is weak at present. Although the impact of these experimental results on the validity of CKM and SM is currently statistically limited, they have attracted much interest in searching for new physics and it has been shown that the deviation can be understood without contradicting the smallness of the SUSY effect on $B \to J/\psi K_S$ in the minimal supersymmetric standard model (MSSM) [106, 107].

Another experimental result is of the time dependent CP asymmetry $S_{\eta' K_S}$ in $B \to \eta' K_S$

$$
\begin{aligned}
S_{\eta' K_S} &= 0.02 \pm 0.34 \pm 0.03 \quad \text{BaBar} \\
&= 0.43 \pm 0.27 \pm 0.05 \quad \text{Belle}
\end{aligned}
\tag{24}
$$

which deviates sizably from the SM expectation. Although both the asymmetries $S_{\phi K}$ and $S_{\eta' K}$ are smaller than the SM value, $S_{\phi K}$ is negative and $S_{\eta' K}$ is positive, which implies that new CP violating physics affects $B \to \phi K_S$ in a dramatic way but gives $B \to \eta' K_S$ a relatively small effect. Because the quark subprocess $b \to s\bar{s}s$ contributes to both $B \to \phi K_S$ and $B \to \eta' K_S$ decays one should simultaneously explain the experimental data in a model with the same parameters. It has been done in Ref. [108] in a model- independent way in the supersymmetric (SUSY) framework. In Ref. [108] the analysis is carried out using the naive factorization to calculate hadronic matrix elements of operators and the neutral Higgs boson (NHB) contributions are not included. As we have shown in a letter [107] that both the branching ratio (Br) and CP asymmetry are significantly dependent of the α_s corrections of hadronic matrix elements and NHB contributions are important in MSSM with middle and large $\tan\beta$ (say, > 8). In this paper [109] we study the $B \to \phi K_S$ and $B \to \eta' K_S$ decays in MSSM by calculating hadronic matrix elements of operators with QCD factorization approach and including neutral Higgs boson (NHB) contributions. We calculate the Wilson coefficients of operators including the new operators which are induced by NHB penguins at LO using the MIA with double insertions. We calculate the α_s order hadronic matrix elements of the new operators for $B \to \phi K_s$ and $B \to \eta' K_s$. We analyze constraints from relevant experimental data among which the new CDF bound of $B_s \to \mu^+\mu^-$ is used. It is shown that the recent experimental results on the time-dependent CP asymmetries in $B \to \phi K_S$ and $B \to \eta' K_S$, $S_{\phi K}$ is negative and $S_{\eta' K}$ is positive but smaller than 0.7, which can not be explained in SM, can be explained in MSSM if there are flavor non-diagonal squark mass matrix elements of 2nd and 3rd generations whose size satisfies all relevant constraints from known experiments ($B \to X_S\gamma, B_s \to \mu^+\mu^-, B \to X_s\mu^+\mu^-, B \to X_s g, \Delta M_s$, etc.). In particular, we find that one can explain the experimental results with a flavor non-diagonal mass insertion of chirality LL or LR when α_s corrections of hadronic matrix elements of operators are included, in contrast with the claim in the literature. At the same time, the branching ratios for the two decays can also be in agreement with experimental measurements.

C. B decays in model III and TC models

In Ref.[110] the authors made a systematic study for the new physics contributions to the charmless two-body hadronic decays of B and B_s meson induced by the neutral- and charged-Higgs penguin diagrams in model III: the third type of the two-Higgs-doublet models. They found that the new physics enhancements to the branching ratios of those penguin dominated B meson decay modes can be significant, about 30% to 50%, but the corresponding new physics corrections to CP violating asymmetries are generally small. Also in model III, the penguin contributions to

the decays of $b \to sg$, $B \to X_s\gamma$, $X_s\gamma\gamma$, $V\gamma$ ($V = K^*, \rho, \omega$), l^+l^- and $B^0 - \bar{B}^0$ mixing have been studied, for example, in Ref.[111]

In the framework of various technicolor models, many B meson decay modes, such as the radiative decays $B \to X_s\gamma$ and $B \to X_s\gamma\gamma$, the semileptonic decays $B \to K^{(*)}l^+l^-$ and $B \to X_s\nu\bar{\nu}$, the charmless two-body hadronic decays $B \to h_1h_2$, have been calculated recently [112]. Strong constraints on the parameter space of technicolor models have been found from these studies.

X. SUMMARY

In this B physics workshop, a number of interesting topic of B physics is discussed. The heavy quark effective theory and heavy quark effective field theory are intensively discussed and compared. The two approaches for hadronic B decays, PQCD and QCD factorization are also discussed and compared. A lot of questions are raised. The hadronic wave function including its twist 3 part are derived. The contribution of new physics especially super symmetry are discussed. Most scientists and students benefit a lot from the lectures and discussions. For a summary, Prof. Yue-Liang Wu writes In his transparencies:

- BOTTOM IS BEAUTY ! $b \leftrightarrow \bar{b}$

- BEAUTY FROM TOP ! $b \leftarrow t$

- BEAUTY TO UP ! $b \to u$

- BEAUTY TO CHARM ! $b \to c$

- BEAUTY TO STRANGE ! $b \to s$

- PHYSICS REMAINS AT BOTTOM , LCQCD/HQEFT/QCDF/PQCD IS VERY USEFUL

- MORE NEW PHYSICS BEYOND SM, PHYSICS IS BEAUTY & CHARM

Acknowledgments

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